Further study of the helicity selection rule evading mechanism in η_c, χ_{c0} and h_c decaying to baryon anti-baryon pairs

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We investigate the long distance contribution via charmed hadron loops in the processes η_c, χ_{c0} and h_c decaying to baryon anti-baryon pairs, which are supposed to be highly suppressed by the helicity selection rule as a consequence of the perturbative QCD framework. With an effective Lagrangian method, our estimation result indicates that such hadron loops play an important role in these hadronic decays. It is a further test of the evading mechanism for the helicity selection rule in charmonium baryon-antibaryon decays.

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I. INTRODUCTION

The complexity of quantum chromodynamics (QCD) remains unsolved in many questions in the charmonium mass region. This is the regime that numerous contradictory results between perturbative calculations and experimental observations were found. In particular, the pQCD expected helicity selection rule [1–3] has been found violated in many exclusive decay processes. According to this selection rule, some charmonium decay channels are supposed to be highly suppressed, such as $J/\psi \to VP$, $\eta_c \to VV$ and $\chi_{c1} \to VV$, where V and P denote vector and pseudoscalar meson respectively. However, all these decays are found to be rather important in experiments [4]. In Ref. [5], Feldmann and Kroll argued that charmonium mesonic decays could be classified into two catalogues. One is controlled by pQCD with conserved hadronic helicity while the other is characterized by the helicity-selection-rule violation. It was proposed that a soft mechanism through a light-quark Fock component would account for the violation of the hadronic helicity conservation in exclusive decays. Other theoretical solutions for the underlying dynamics have also been explored in the literature [5–10], and detailed review of some of the related questions in charmonium exclusive decays can be found in Refs. [1–3, 11–13].

In this work we are going to investigate another three processes, i.e., $\eta_c, \chi_{c0}, h_c \to Y\bar{Y}$, where $Y\bar{Y}$ represent the $J^P = 1/2^+$ octet baryon-antibaryon pairs, i.e. $p\bar{p}$, $\Lambda\bar{\Lambda}$, $\Sigma\bar{\Sigma}$, and $\Xi\bar{\Xi}$. These channels are also supposed to be highly suppressed according to the helicity selection rule [14]. However, it seems that the available experimental data do not support such expectations at all [4]. Some attempts have been made in order to understand this contradiction [15–20]. For instance, in the processes of η_c and $\chi_{cJ}(J = 0, 1, 2)$ decaying into $p\bar{p}$, a quark-diquark model for proton is introduced as a mechanism for evading the helicity selection rule. However, the obtained branching ratios for η_c and χ_{cJ} can not simultaneously agree with the experimental data [15, 16]. Constituent quark mass corrections were also introduced to account for the processes $\eta_c, \chi_{c0}, h_c \to p\bar{p}$ [17, 18], where it turns out that the obtained branching ratio for $\eta_c \to p\bar{p}$ is still much smaller than the experimental data. There are also some other proposed mechanisms for understanding these processes, such as the mixing between charmonium state and glueball [19], and the quark pair creation model [20]. Basically, it is believed that non-perturbative mechanisms should be relevant in this scenario.

In Refs. [8, 9], it was proposed that intermediate meson loop (IML) transitions can serve as a soft mechanism in charmonium decays. Such a long-distance interaction can evade the Okubo-Zweig-Iizuka (OZI) rule and result in violation of the pQCD helicity selection rule. In the framework of effective Lagrangians for hadron interactions, we can then quantitatively study the evasion of the helicity selection rule in various processes. As a further check of this mechanism, we are going to investigate the role of the charmed hadron loops in $\eta_c, \chi_{c0}, h_c \to Y\bar{Y}$ in this work. Additional evidence for such an underlying mechanism should allow us to gain more insights into the QCD strong interaction properties in the intermediate energy region.

The rest of this article is arranged as follows: In Sec. II, we will depict the effective Lagrangian method, and some relevant formulas will be given. In Sec. III, we will present the numerical results and discussions. The conclusions will be summarized in Sec. IV.

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II. LONG-DISTANCE CONTRIBUTION VIA CHARMED HADRON LOOPS

In Figs. 1, 2, 3 and 4, the intermediate charmed hadron loops which serve as a long-distance soft mechanism are illustrated by triangle diagrams. Since one charmed baryon will be present in the loop, these diagrams are somewhat different from the intermediate meson loops studied in Refs. [8, 9]. We will consider the exchanges of the ground state $J^P = 1/2^+$ charmed baryons that belong to SU(4) multiplets [4, 21]. There are several points that should be clarified for Figs. 1, 2, 3 and 4:

1) $J_{c\bar{c}}$ represents the decaying $c\bar{c}$ meson, i.e. η_c, χ_{c0}, or h_c.

2) $\mathbb{D}_{(s)}$ and $\bar{\mathbb{D}}_{(s)}$ only denote the flavor contents but do not contain the spin quantum numbers, and they will depend on the decaying mesons.

3) Different intermediate charmed hadrons will appear in the loops for $\eta_c, \chi_{c0}, h_c \to Y\bar{Y}$. For the η_c decay, there are three situations for $\mathbb{D}_{(s)}$ and $\bar{\mathbb{D}}_{(s)}$: (a) $D_{(s)}\bar{D}^*_{(s)}$, (b) $D^*_{(s)}\bar{D}_{(s)}$, (c) $D^*_{(s)}\bar{D}^*_{(s)}$. For the χ_{c0} decay there are two: (a) $D_{(s)}\bar{D}_{(s)}$, (b) $D^*_{(s)}\bar{D}^*_{(s)}$. For the h_c decay there are three: (a) $D_{(s)}\bar{D}^*_{(s)}$, (b) $D^*_{(s)}\bar{D}_{(s)}$, (c) $D^*_{(s)}\bar{D}^*_{(s)}$. This is due to the adopted effective Lagrangians based on heavy quark symmetry [22, 23].

We list the relevant effective Lagrangians as follows:

$$\mathcal{L}_1 = ig_1 Tr[P^\mu_{c\bar{c}}\bar{H}_{2i}\gamma_\mu \bar{H}_{1i}] + h.c., \tag{1}$$

$$\mathcal{L}_2 = ig_2 Tr[R_{c\bar{c}}\bar{H}_{2i}\gamma^\mu \overleftrightarrow{\partial}_\mu \bar{H}_{1i}] + h.c., \tag{2}$$

where the spin multiplets for these four P-wave and two S-wave charmonium states are expressed as

$$P^\mu_{c\bar{c}} = \left(\frac{1+\not{v}}{2}\right)\left(\chi^{\mu\alpha}_{c2}\gamma_\alpha + \frac{1}{\sqrt{2}}\epsilon^{\mu\nu\alpha\beta}v_\alpha\gamma_\beta\chi_{c1\nu} + \frac{1}{\sqrt{3}}(\gamma^\mu - v^\mu)\chi_{c0} + h^\mu_c\gamma_5\right)\left(\frac{1-\not{v}}{2}\right), \tag{3}$$

$$R_{c\bar{c}} = \left(\frac{1+\not{v}}{2}\right)(\psi^\mu\gamma_\mu - \eta_c\gamma_5)\left(\frac{1-\not{v}}{2}\right). \tag{4}$$

The charmed and anti-charmed meson triplet read

$$H_{1i} = \left(\frac{1+\not{v}}{2}\right)[\mathcal{D}^{*\mu}_i\gamma_\mu - \mathcal{D}_i\gamma_5], \tag{5}$$

$$H_{2i} = [\bar{\mathcal{D}}^{*\mu}_i\gamma_\mu - \bar{\mathcal{D}}_i\gamma_5]\left(\frac{1-\not{v}}{2}\right), \tag{6}$$

where $\mathcal{D}$ and $\mathcal{D}^*$ denote the pseudoscalar and vector charmed meson fields respectively, i.e. $\mathcal{D}^{(*)} = (D^{0(*)}, D^{+(*)}, D_s^{+(*)})$. For the meson-baryon interaction Lagrangians, we follow the forms that were adopted in Refs. [24–27]:

$$\mathcal{L}_{Y_c\mathcal{D}Y} = ig_{Y_c\mathcal{D}Y}\bar{Y}_c\gamma_5 Y\mathcal{D} + h.c., \tag{7}$$

$$\mathcal{L}_{Y_c\mathcal{D}^*Y} = g_{Y_c\mathcal{D}^*Y}\bar{Y}_c\left(\gamma^\mu\mathcal{D}^*_\mu + \frac{\kappa_{Y_c\mathcal{D}^*Y}}{2m_N}\sigma^{\mu\nu}\partial_\mu\mathcal{D}^*_\nu\right)Y + h.c., \tag{8}$$

where Y_c, $\mathcal{D}^{(*)}$, and Y denote the charmed baryon, charmed meson, and the corresponding nucleon or hyperon, respectively. The relevant coupling constants will be discussed later.

With the above effective Lagrangians, we can now calculate the transition amplitudes illustrated in Figs. 1, 2, 3 and 4. For these diagrams, we take the convention of the momenta as $J_{c\bar{c}}(p) \to \mathbb{D}(q_1)\bar{\mathbb{D}}(q_2)[Y_c(q)] \to \bar{Y}(p_1)Y(p_2)$.

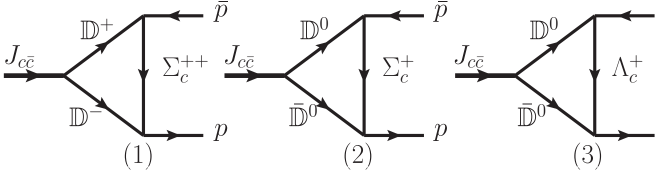

FIG. 1: Charmed hadron loop diagrams that describe the long-distance transitions in $J_{c\bar{c}} \to p\bar{p}$.

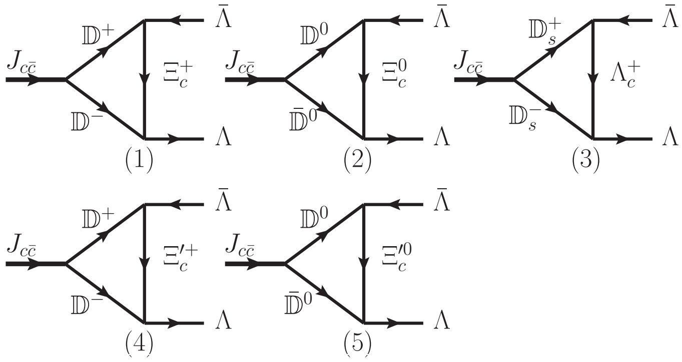

FIG. 2: Charmed hadron loop diagrams that describe the long-distance transitions in $J_{c\bar{c}} \to \Lambda\bar{\Lambda}$.

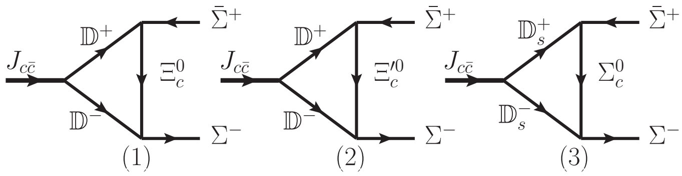

FIG. 3: Charmed hadron loop diagrams that describe the long-distance transitions in $J_{c\bar{c}} \to \Sigma^-\bar{\Sigma}^+$.

A. $\eta_c \to Y\bar{Y}$

We first consider $\eta_c \to Y\bar{Y}$, where three different amplitudes corresponding to those three intermediate states will contribute, i.e. $\mathcal{M}_a$, $\mathcal{M}_b$, and $\mathcal{M}_c$. These amplitudes can be written down explicitly as follows:

$$\begin{aligned}
\mathcal{M}_a &= 2g_{\eta_c\mathcal{D}\mathcal{D}^*}g_{Y_c\mathcal{D}Y}g_{Y_c\mathcal{D}^*Y}\int\frac{d^4q}{(2\pi)^4}(q_{2\lambda}-q_{1\lambda})\left(-g^{\lambda\mu}+\frac{q_2^\lambda q_2^\mu}{m_{\mathcal{D}^*}^2}\right)\\
&\times\ \bar{u}(p_2)\left(\gamma_\mu+i\frac{\kappa_{Y_c\mathcal{D}^*Y}}{2m_N}\sigma_{\mu\nu}q_2^\nu\right)(\not{q}+m_{Y_c})\gamma_5 v(p_1)\\
&\times\ \frac{1}{q^2-m_{Y_c}^2}\frac{1}{q_1^2-m_{\mathcal{D}}^2}\frac{1}{q_2^2-m_{\mathcal{D}^*}^2}\mathcal{F}(q^2),
\end{aligned}\tag{9}$$

$$\begin{aligned}
\mathcal{M}_b &= 2g_{\eta_c\mathcal{D}\mathcal{D}^*}g_{Y_c\mathcal{D}Y}g_{Y_c\mathcal{D}^*Y}\int\frac{d^4q}{(2\pi)^4}(q_{2\lambda}-q_{1\lambda})\left(-g^{\lambda\mu}+\frac{q_1^\lambda q_1^\mu}{m_{\mathcal{D}^*}^2}\right)\\
&\times\ \bar{u}(p_2)\gamma_5(\not{q}+m_{Y_c})\left(\gamma_\mu+i\frac{\kappa_{Y_c\mathcal{D}^*Y}}{2m_N}\sigma_{\mu\nu}q_1^\nu\right)v(p_1)\\
&\times\ \frac{1}{q^2-m_{Y_c}^2}\frac{1}{q_1^2-m_{\mathcal{D}^*}^2}\frac{1}{q_2^2-m_{\mathcal{D}}^2}\mathcal{F}(q^2),
\end{aligned}\tag{10}$$

$$\begin{aligned}
\mathcal{M}_c &= -2ig_{\eta_c\mathcal{D}^*\mathcal{D}^*}g_{Y_c\mathcal{D}^*Y}^2\int\frac{d^4q}{(2\pi)^4}\epsilon^{\mu\nu\lambda\tau}p_\nu(q_{2\mu}-q_{1\mu})\\
&\times\ \bar{u}(p_2)\left(\gamma_\tau+i\frac{\kappa_{Y_c\mathcal{D}^*Y}}{2m_N}\sigma_{\tau\xi}q_2^\xi\right)(\not{q}+m_{Y_c})\left(\gamma_\lambda+i\frac{\kappa_{Y_c\mathcal{D}^*Y}}{2m_N}\sigma_{\lambda\sigma}q_1^\sigma\right)v(p_1)
\end{aligned}$$

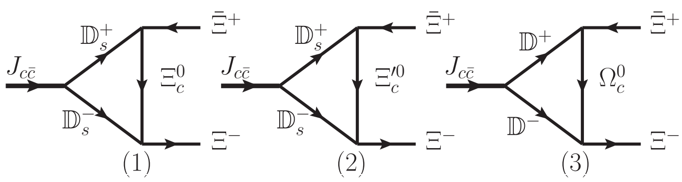

FIG. 4: Charmed hadron loop diagrams that describe the long-distance transitions in $J_{c\bar{c}} \to \Xi^-\bar{\Xi}^+$.

$$\times \frac{1}{q^2 - m_{Y_c}^2} \frac{1}{q_1^2 - m_{\mathcal{D}^*}^2} \frac{1}{q_2^2 - m_{\mathcal{D}^*}^2} \mathcal{F}(q^2). \tag{11}$$

Since the masses of η_c, χ_{c0}, and h_c are below the threshold of $D^{(*)}\bar{D}^{(*)}$, intermediate mesons $D^{(*)}$ and $\bar{D}^{(*)}$ can not be on-shell simultaneously. We thus phenomenologically introduce a form factor $\mathcal{F}(q^2)$ as has been done in Refs. [9, 22, 28] to take into account the off-shell effects,

$$\mathcal{F}(q^2) = \prod_i \left(\frac{m_i^2 - \Lambda_i^2}{q_i^2 - \Lambda_i^2} \right), \tag{12}$$

where $q_i = q,\ q_1,\ q_2$. The cut-off energy is chosen as $\Lambda_i = m_i + \alpha \Lambda_{QCD}$, $\Lambda_{QCD} = 0.22$ GeV and m_i is the mass of the corresponding exchanged particle. The form factor is also necessary for killing the divergence of the loop integrals, although it will give rise to model-dependent aspects of the calculations. Qualitatively, this type of form factor will converge the integrals faster than a dipole form factor and can be well combined with a vertex coupling constant as a phenomenological account of the non-local coupling form factor. Some further discussions will be given later in the numerical analysis.

In this work, we do not include contributions from exchanging other charmed hadrons with higher spin or orbital excitations. These resonances have relatively larger masses and their couplings so far are unknown. It should be reasonable to only consider the lowest partial wave states based on the argument of locally broken down quark-hadron duality [29].

B. $\chi_{c0} \to Y\bar{Y}$

Two charmed hadron loops will contribute to $\chi_{c0} \to Y\bar{Y}$, and the corresponding amplitudes are

$$\begin{aligned}
\mathcal{M}_a &= -i g_{\chi_{c0}\mathcal{D}\mathcal{D}} g_{Y_c\mathcal{D}Y}^2 \int \frac{d^4q}{(2\pi)^4} \bar{u}(p_2)\gamma_5(\not{q} + m_{Y_c})\gamma_5 v(p_1) \\
&\times \frac{1}{q^2 - m_{Y_c}^2} \frac{1}{q_1^2 - m_{\mathcal{D}}^2} \frac{1}{q_2^2 - m_{\mathcal{D}}^2} \mathcal{F}(q^2),
\end{aligned} \tag{13}$$

$$\begin{aligned}
\mathcal{M}_b &= i g_{\chi_{c0}\mathcal{D}^*\mathcal{D}^*} g_{Y_c\mathcal{D}^*Y}^2 \int \frac{d^4q}{(2\pi)^4} g_{\mu\nu} \left(-g^{\mu\rho} + \frac{q_1^\mu q_1^\rho}{m_{\mathcal{D}^*}^2} \right) \left(-g^{\nu\alpha} + \frac{q_2^\nu q_2^\alpha}{m_{\mathcal{D}^*}^2} \right) \\
&\times \bar{u}(p_2) \left(\gamma_\alpha + i \frac{\kappa_{Y_c\mathcal{D}^*Y}}{2m_N} \sigma_{\alpha\beta} q_2^\beta \right) (\not{q} + m_{Y_c}) \left(\gamma_\rho + i \frac{\kappa_{Y_c\mathcal{D}^*Y}}{2m_N} \sigma_{\rho\tau} q_1^\tau \right) v(p_1) \\
&\times \frac{1}{q^2 - m_{Y_c}^2} \frac{1}{q_1^2 - m_{\mathcal{D}^*}^2} \frac{1}{q_2^2 - m_{\mathcal{D}^*}^2} \mathcal{F}(q^2),
\end{aligned} \tag{14}$$

where the form factor has the same form as that in $\eta_c \to Y\bar{Y}$.

C. $h_c \to Y\bar{Y}$

Similarly, the amplitudes for $h_c \to Y\bar{Y}$ from those three contributing loops can be written down as follows,

$$\begin{aligned}
\mathcal{M}_a &= g_{h_c\mathcal{D}\mathcal{D}^*} g_{Y_c\mathcal{D}Y} g_{Y_c\mathcal{D}^*Y} \epsilon_\eta(p) \int \frac{d^4q}{(2\pi)^4} \left(-g^{\lambda\mu} + \frac{q_2^\lambda q_2^\mu}{m_{\mathcal{D}^*}^2} \right) \\
&\times \bar{u}(p_2) \left(\gamma_\mu + i \frac{\kappa_{Y_c\mathcal{D}^*Y}}{2m_N} \sigma_{\mu\nu} q_2^\nu \right) (\not{q} + m_{Y_c})\gamma_5 v(p_1) \\
&\times \frac{1}{q^2 - m_{Y_c}^2} \frac{1}{q_1^2 - m_{\mathcal{D}}^2} \frac{1}{q_2^2 - m_{\mathcal{D}^*}^2} \mathcal{F}(q^2),
\end{aligned} \tag{15}$$

$$\begin{aligned}
\mathcal{M}_b &= g_{h_c\mathcal{D}\mathcal{D}^*} g_{Y_c\mathcal{D}Y} g_{Y_c\mathcal{D}^*Y} \epsilon_\eta(p) \int \frac{d^4q}{(2\pi)^4} \left(-g^{\lambda\mu} + \frac{q_1^\lambda q_1^\mu}{m_{\mathcal{D}^*}^2} \right) \\
&\times \bar{u}(p_2)\gamma_5(\not{q} + m_{Y_c}) \left(\gamma_\mu + i \frac{\kappa_{Y_c\mathcal{D}^*Y}}{2m_N} \sigma_{\mu\nu} q_1^\nu \right) v(p_1)
\end{aligned}$$

$$\times \frac{1}{q^2 - m_{Y_c}^2} \frac{1}{q_1^2 - m_{\mathcal{D}^*}^2} \frac{1}{q_2^2 - m_{\mathcal{D}}^2} \mathcal{F}(q^2), \tag{16}$$

$$\mathcal{M}_c = g_{h_c\mathcal{D}^*\mathcal{D}^*} g_{Y_c\mathcal{D}^*Y}^2 \epsilon^\eta(p) \int \frac{d^4q}{(2\pi)^4} \epsilon_{\rho\eta\alpha\beta} p^\rho \left(-g^{\alpha\lambda} + \frac{q_1^\alpha q_1^\lambda}{m_{\mathcal{D}^*}^2}\right) \left(-g^{\beta\tau} + \frac{q_2^\beta q_2^\tau}{m_{\mathcal{D}^*}^2}\right)$$
$$\times \bar{u}(p_2) \left(\gamma_\tau + i\frac{\kappa_{Y_c\mathcal{D}^*Y}}{2m_N} \sigma_{\tau\xi} q_2^\xi\right) (\not{q} + m_{Y_c}) \left(\gamma_\lambda + i\frac{\kappa_{Y_c\mathcal{D}^*Y}}{2m_N} \sigma_{\lambda\sigma} q_1^\sigma\right) v(p_1)$$
$$\times \frac{1}{q^2 - m_{Y_c}^2} \frac{1}{q_1^2 - m_{\mathcal{D}^*}^2} \frac{1}{q_2^2 - m_{\mathcal{D}^*}^2} \mathcal{F}(q^2), \tag{17}$$

where $\epsilon_\eta(p)$ is the polarization vector for h_c, and again the form factor has the same form as that in $\eta_c \to Y\bar{Y}$.

III. NUMERICAL RESULTS AND DISCUSSION

Proceeding to the numerical results, we first discuss the determination of coupling constants. For the couplings of the charmonium states to charmed mesons, the expansion of the effective Lagrangians $\mathcal{L}_1$ and $\mathcal{L}_2$ gives the following relations in the heavy quark limit:

$$g_{\eta_c\mathcal{D}\mathcal{D}^*} = 2g_2\sqrt{m_{\eta_c} m_{\mathcal{D}} m_{\mathcal{D}^*}}, \; g_{\eta_c\mathcal{D}^*\mathcal{D}^*} = 2g_2 \frac{m_{\mathcal{D}^*}}{\sqrt{m_{\eta_c}}},$$
$$g_{\chi_{c0}\mathcal{D}\mathcal{D}} = -2\sqrt{3} g_1 m_{\mathcal{D}} \sqrt{m_{\chi_{c0}}}, \; g_{\chi_{c0}\mathcal{D}^*\mathcal{D}^*} = -\frac{2}{\sqrt{3}} g_1 m_{\mathcal{D}^*} \sqrt{m_{\chi_{c0}}},$$
$$g_{h_c\mathcal{D}\mathcal{D}^*} = -2g_1\sqrt{m_{h_c} m_{\mathcal{D}} m_{\mathcal{D}^*}}, \; g_{h_c\mathcal{D}^*\mathcal{D}^*} = 2g_1 \frac{m_{\mathcal{D}^*}}{\sqrt{m_{h_c}}},$$
$$g_1 = -\sqrt{\frac{m_{\chi_{c0}}}{3}} \frac{1}{f_{\chi_{c0}}}, \; g_2 = \frac{\sqrt{m_\psi}}{2m_{\mathcal{D}} f_\psi}, \tag{18}$$

where $f_{\chi_{c0}}$ and f_ψ are the decay constants of χ_{c0} and J/ψ, respectively. $f_\psi \simeq 0.41$ GeV, which is determined by the experimental data [4]. And $f_{\chi_{c0}}$ can be approximately determined by the QCD sum rule approach, i.e. $f_{\chi_{c0}} \simeq 0.51$ GeV [30].

There is no much information on the couplings of a charmed baryon to a charmed meson and light baryon. If considering SU(4) flavor symmetry, it will relate different couplings with each other. For the considered $J^P = 1/2^+$ charmed baryons, we would expect the following relations [31]:

$$[\bar{\Sigma}_c^{--} D^+ p] = [\bar{\Sigma}_c^{--} D_s^+ \Sigma^+] = -\sqrt{2}[\bar{\Xi}_c^{\prime 0} D^+ \Sigma^-] = -\sqrt{2}[\bar{\Xi}_c^{\prime -} D_s^+ \Xi^0]$$
$$= -\frac{2}{\sqrt{3}}[\bar{\Xi}_c^{\prime -} D^+ \Lambda] = -[\bar{\Omega}_c^0 D^+ \Xi^-] = -\sqrt{2}[\bar{p} K^+ \Sigma^0],$$
$$[\bar{\Lambda}_c^- D^+ n] = [\bar{\Xi}_c^0 D^+ \Sigma^-] = -[\bar{\Xi}_c^- D_s^+ \Xi^0] = -\sqrt{6}[\bar{\Xi}_c^- D^+ \Lambda]$$
$$= \sqrt{\frac{3}{2}}[\bar{\Lambda}_c^- D_s^+ \Lambda] = -[\bar{p} K^+ \Lambda], \tag{19}$$

where the square bracket "$[\cdots]$" denotes the coupling constant for the corresponding vertex, and the pseudoscalar meson in the bracket can be replaced by the corresponding vector meson. These formulas relate the charmed-baryon-meson couplings with the strange-baryon-meson coupling. For instance, the commonly adopted value for ΛKN coupling is $g_{\Lambda KN} = -13.2$ [32–34], from which we obtain $g_{\Lambda_c DN} = -g_{\Lambda KN} = 13.2$. Such a relation may contain rather large uncertainties as we can see that the QCD sum rule approach suggests a smaller value $|g_{\Lambda_c DN}| = 6.7 \pm 2.1$ [35]. Experimental data from the charmed meson photoproduction and charmonium absorption by nucleons may offer some constraints on these couplings. Unfortunately, the extractions of the couplings still depend on the adopted theoretical models [24, 25, 27]. As a tentative solution for this, we empirically retain the SU(4) symmetry relations in the numerical calculation. The relevant strange-baryon-meson couplings are taken from the Nigmegen potential model as follows [32, 33]:

$$g_{\Lambda K^* N} = -4.26, \; \kappa_{\Lambda K^* N} = 2.16,$$
$$g_{\Sigma K^* N} = -2.5, \; \kappa_{\Sigma K^* N} = -0.22, \tag{20}$$

BR(in units of 10^{-4})	$p\bar{p}$	$\Lambda\bar{\Lambda}$	$\Sigma^-\bar{\Sigma}^+$	$\Xi^-\bar{\Xi}^+$
Hadron loop	$9.0 \sim 17.0$	$6.3 \sim 12.5$	$5.05 \sim 10.0$	$4.82 \sim 9.56$
Exp.	13 ± 4	10.4 ± 3.1	-	-

TABLE I: Branching ratios for $\eta_c \to Y\bar{Y}$ predicted by the intermediate charmed hadron loop transitions in the range $\alpha = 0.47 \sim 0.53$ which corresponds to the measured lower and upper bound of BR($\eta_c \to p\bar{p}$) . The available experimental data are taken from Ref. [4], and the dashes mean that the data are unavailable.

and

$$g_{\Lambda KN} = -13.2,\ g_{\Sigma KN} = 3.9. \tag{21}$$

The form factor parameter α generally cannot be determined from the first principle. It will depend on a particular process and its value is order of unity. There might exist some differences between the values of the form factor parameter for $\eta_c \to p\bar{p}$ and $\chi_{c0} \to p\bar{p}$ since η_c and χ_{c0} belong to different spin multiplets. It is natural to anticipate that the counter term structures would be different for loops involving different spin multiplets. Similar feature was also found in Ref. [28] and other studies [9]. As emphasized in the literature, in order to reduce the uncertainties arising from the form factor, one first has to rely on the experimental data to determine the form factor parameter, and then apply it to other flavor-symmetry-related processes to make predictions which can be examined by further experimental data. This empirical treatment will also be useful for a further control of the uncertainties arising from the SU(4) flavor symmetry breaking. Namely, with the charmed-baryon-meson couplings fixed by the SU(4) relations, we allow the experimental data for $\eta_c \to p\bar{p}$ to determine the range of α as a compensation of the uncertainties from the couplings. This can be regarded reasonable since a physical coupling form factor will generally be correlated with these two aspects. Reliability of this treatment can be tested by the prediction for the $\eta_c \to \Lambda\bar{\Lambda}$ branching ratio, which turns out to be consistent with the experimental data [4].

At present, the branching ratios for $\eta_c \to p\bar{p}$ and $\chi_{c0} \to p\bar{p}$ have been measured by experiment [4]. Signals of h_c were also found in $p\bar{p}$ annihilations [36, 37], which is a hint that $h_c \to p\bar{p}$ could be an important channel in the h_c decays. Thus, we will use the measured branching ratios BR($\eta_c \to p\bar{p}$) and BR($\chi_{c0} \to p\bar{p}$) to extract the form factor parameters for the S-wave state η_c and P-wave state χ_{c0}, respectively. Considering that h_c belongs to the same spin multiplet as χ_{c0}, we conjecture that they may share the same intrinsic dynamics in their decays into $p\bar{p}$. We will then adopt the same form factor parameter for h_c as that extracted from $\chi_{c0} \to p\bar{p}$.

We use the software package LoopTools to calculate the loop integrals [38]. The results are displayed in Tables I, II, and III. The experimental data for $\eta_c \to p\bar{p}$ and $\chi_{c0} \to p\bar{p}$ are adopted for the determination of the form factor parameter α with the range of the uncertainties. In Table I, the predicted branching ratio, BR($\eta_c \to \Lambda\bar{\Lambda}$) = $(6.3 \sim 12.5) \times 10^{-4}$, is consistent with the data BR^{exp}($\eta_c \to \Lambda\bar{\Lambda}$) = $(10.4 \pm 3.1) \times 10^{-4}$ [4], which is a sign for the parameter under control. The experimental data for $\eta_c \to \Sigma\bar{\Sigma}$ and $\Xi\bar{\Xi}$ are unavailable. Our calculations suggest that these two branching ratios are compatible with that for $\eta_c \to \Lambda\bar{\Lambda}$. This expectation can be examined by the BESIII experiment.

The experimental data for $\chi_{c0} \to p\bar{p}$ and $\Lambda\bar{\Lambda}$ will allow a further check of the model and its parameter space in the spin-1 multiplets. In Table II, the calculation results are given by the form factor parameter within a range, i.e. $\alpha = 0.23 \sim 0.24$, which corresponds to the experimental uncertainties of $\chi_{c0} \to p\bar{p}$ [4]. This value range is different from that for the η_c decays with $\alpha = 0.47 \sim 0.53$. This can be understood as a consequence that χ_{c0} and η_c belong to different spin multiplets. We shall show later the sensitivities of the calculation results to the form factor parameter later. For the χ_{c0} decays, the data from CLEO suggest relatively larger branching ratios for $\chi_{c0} \to \Lambda\bar{\Lambda}$ compared with that for $\chi_{c0} \to p\bar{p}$. Also, branching ratios for $\chi_{c0} \to \Sigma\bar{\Sigma}$ and $\Xi\bar{\Xi}$ are sizeable. In contrast, with the same form factor parameter α, we find relatively smaller branching ratios for $\chi_{c0} \to \Lambda\bar{\Lambda}$, $\Sigma\bar{\Sigma}$ and $\Xi\bar{\Xi}$. It could be a sign that the SU(4) flavor symmetry is badly broken. Namely, Eq. (19) may be too rough, and can only provide an estimate of magnitude orders for $\chi_{c0} \to \Sigma\bar{\Sigma}$ and $\Xi\bar{\Xi}$. In this sense, the calculation results for $\chi_{c0} \to \Lambda\bar{\Lambda}$, $\Sigma\bar{\Sigma}$ and $\Xi\bar{\Xi}$, though turn out to be smaller than the experimental data, can be regarded as reasonable.

Although the present model uncertainties do not allow us to conclude the magnitudes of the hadron loop contributions, the pattern predicted by this mechanism still suggests an important role played by the hadron loops in the explanation of the helicity-selection-rule violations in these exclusive decays. Meanwhile, we stress that it would be essential to have improved experimental measurements in order to gain better insights into the transition mechanism.

BR(in units of 10^{-4})	$p\bar{p}$	$\Lambda\bar{\Lambda}$	$\Sigma^-\bar{\Sigma}^+$	$\Xi^-\bar{\Xi}^+$
Hadron loop	$1.96 \sim 2.34$	$1.19 \sim 1.51$	$0.55 \sim 0.69$	$0.52 \sim 0.66$
Exp. [4]	2.15 ± 0.19	4.4 ± 1.5	-	< 10.3
Exp. [39]	2.25 ± 0.27	4.7 ± 1.6	3.25 ± 1.14	5.14 ± 1.25

TABLE II: Branching ratios for $\chi_{c0} \to Y\bar{Y}$ predicted by the intermediate charmed hadron loop transitions in the range $\alpha = 0.23 \sim 0.24$ which corresponds to the measured lower and upper bound of BR($\chi_{c0} \to p\bar{p}$) [4].

BR(in units of 10^{-4})	$p\bar{p}$	$\Lambda\bar{\Lambda}$	$\Sigma^-\bar{\Sigma}^+$	$\Xi^-\bar{\Xi}^+$
Hadron loop	$15.2 \sim 19.3$	$5.88 \sim 7.47$	$4.56 \sim 5.80$	$5.57 \sim 7.08$
Exp.	-	-	-	-

TABLE III: Branching ratios for $h_c \to Y\bar{Y}$ predicted by the intermediate charmed hadron loop transitions. The α range for $h_c \to Y\bar{Y}$ is taken the same as that for $\chi_{c0} \to Y\bar{Y}$, i.e. $\alpha = 0.23 \sim 0.24$. The dashes mean that the data are unavailable. We take the width of h_c as $\Gamma(h_c) = 0.73$ MeV, which is the central value measured by BESIII recently [40].

Experimental data for $h_c \to Y\bar{Y}$ so far are unavailable. Thus, our predictions are based on the assumption that h_c shares the same form factor parameter as χ_{c0} since they belong to the same spin multiplet. In Table III we list the branching ratios given by $\alpha = 0.23 \sim 0.24$, which is the same as adopted in $\chi_{c0} \to p\bar{p}$. In comparison with the theoretical calculations of Refs. [18, 41, 42], our prediction of the branching ratio of $h_c \to p\bar{p}$ seems to be larger. In fact, the theoretical predictions in the literature also appear to be quite different from each other. In particular, some of those results strongly depend on the evaluation of BR($h_c \to J/\psi\pi^0$) in the combined cross sections for $p\bar{p} \to h_c \to J/\psi\pi^0$ from the E760 data [18, 36, 41]. Another reason for the discrepancies among these theoretical estimates may be due to different intrinsic mechanisms adopted for the explanation of the helicity-selection-rule violation. We expect that the future precise measurement of $h_c \to Y\bar{Y}$ will help disentangle the underlying mechanisms.

In Figs. 5, 6 and 7, we also examine the dependence of the results on the form factor parameter α. The adopted values of α are within a reasonable range and well controlled by the available experimental data. Although some uncertainties will be inevitably introduced by the phenomenological form factor, the branching ratio fractions among the considered channels turn out to be stable and less model-dependent. This feature suggests that the branching ratio fractions are less model-dependent in comparison with the absolute branching ratios. In another word, although the model predictions for the absolute branching ratios are lack of experimental constraints, thus, becomes sensitive to the form factor parameter α, we would expect that the predicted branching ratio fractions among those considered channels are less sensitive to it. As a consequence, any experimental results for BR($h_c \to Y\bar{Y}$) at $\mathcal{O}(10^{-4}) \sim \mathcal{O}(10^{-3})$ would imply the importance of the hadron loop contributions to this helicity-selection-rule violation transition.

In the P-wave charmonium decays the next higher Fock state $c\bar{c}g$, i.e. the so-called color octet, will contribute at the same order as the color singlet $c\bar{c}$ in the framework of perturbative factorization method [5, 43–46]. This scenario may share the same intrinsic physics with those intermediate charmed hadron loop transitions based on quark-hadron duality argument, i.e. a manifestation of the same physics at either quark-gluon level, or hadron level. Some qualitative discussions can be found in Refs. [22, 28, 47] and references therein. We address that since it is not easy to handle these exclusive decays with the perturbative methods considering the relatively low energy scale, the intermediate charmed hadron loop transition would serve as a natural soft mechanism in the numerical exploration.

IV. SUMMARY

In this work, with an effective Lagrangian method based on heavy quark and flavor symmetry, we investigate the role played by the intermediate charmed hadron loops in the processes $\eta_c, \chi_{c0}, h_c \to Y\bar{Y}$, which are supposed to be highly suppressed by the helicity selection rule. The results indicate that the transitions via these kinds of loops as long-distance effects can give significant contributions. This is a further test of the mechanism for the evasion of helicity selection rule that we proposed in Ref. [9]. Although the model bares uncertainties arising from the unknown coupling constants and form factor parameter, the available data have provided a reasonable control on the range of the form factor parameter values. Branching ratios for some unmeasured channels can thus be predicted. Sizeable data samples on charmonium decays accumulated at BEPCII and CLEO-c, and future proton-antiproton annihilation data from Panda, are expected to provide a great opportunity for revealing the underlying mechanisms for charmonium helicity-selection-rule-evading decays.

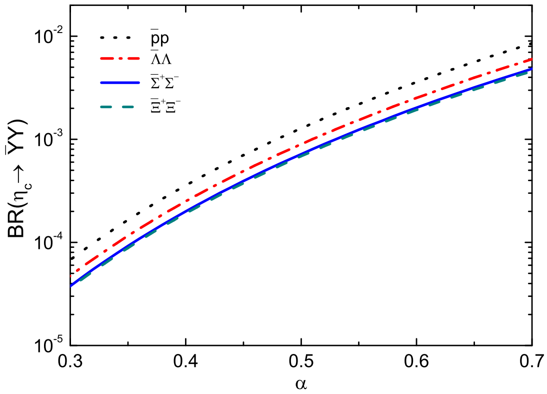

FIG. 5: α-dependence of the calculated branching ratios for $\eta_c \to Y\bar{Y}$.

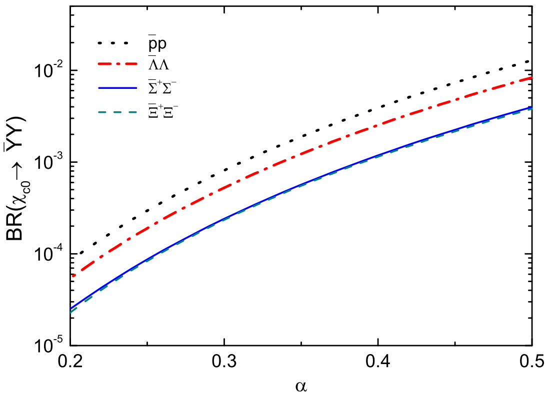

FIG. 6: α-dependence of the calculated branching ratios for $\chi_{c0} \to Y\bar{Y}$.

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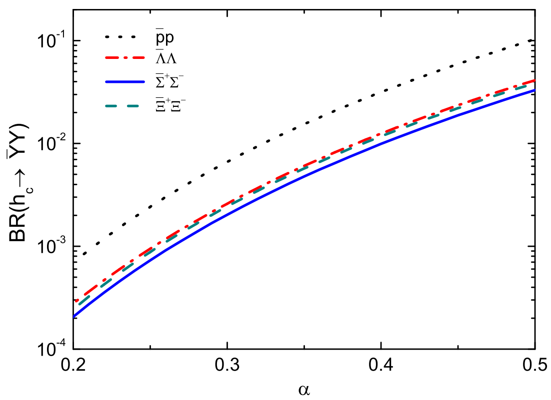

FIG. 7: α-dependence of the calculated branching ratios for $h_c \to Y\bar{Y}$.

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